Collective modes and the far-infrared absorption of the two-dimensional electron gas in a periodic quantizing magnetic field

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Abstract

We investigate the far-infrared (FIR) absorption of a two-dimensional electron gas in a periodically modulated quantizing magnetic field. The magnetic field varies along only one spatial direction and the external time-dependent electric field is linearly polarized along that axis. The mutual Coulomb interaction of the electrons is treated self-consistently in the ground state and in the absorption calculation within the Hartree approximation. The effects of the magnetic material on top of the heterostructure as a grating coupler is included in the time-dependent incident FIR electric field. We show that similar to an electric modulation, the absorption can be directly correlated to the underlying electronic energy bands. In addition, the magnetic modulation leads to absorption spectra with a richer structure due to the quite different static response of the electron density to the modulation.

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I. INTRODUCTION

There is a recent, increasing interest in fabricating and understanding the behavior of semiconductor devices incorporating an electron gas, usually planar, and a nonuniform, or periodic magnetic field (magnetic modulation) of a nanometric length scale. [1–6] The basic methods consist in building metallic patterns on the top of the device, of ferromagnetic or superconducting material, which in the presence of a uniform external magnetic field give a position dependent contribution to the total magnetic field in the plane of the two-dimensional electron gas (2DEG). [7]

Several theoretical papers have analyzed the commensurability oscillations of the resistivity of such systems for weak magnetic fields, i.e. in the classical regime, [8–10] by analogy with the Weiss oscillations, occurring in the presence of a weak electric modulation. For the quantum regime, of strong magnetic fields, magnetic nanostructures like magnetic barriers and magnetic wells have been recently considered. [11,12] For periodic magnetic fields in two spatial directions another recent paper has discussed the energy spectrum equivalent to Hofstadter's butterfly. [13]

So far only very few papers have been dedicated to many-body effects. Wu and Ulloa, [14] have studied the electron density and the collective excitations in a one-dimensional magnetic superlattice with a short period, of the order of the magnetic length corresponding to the average magnetic field. They have found a rapid transition from a two-dimensional to a one-dimensional behavior of the 2DEG, with increasing amplitude of the magnetic modulation. Such a calculation requires the inclusion of the density response of the electron gas to the periodic magnetic field, self-consistently with the energy spectrum. The energy spectrum consists of Landau bands, the degeneracy of the former Landau levels being lifted by the nonuniform magnetic field. Consequently, the electron density becomes nonuniform and reacts to the induced electric field which tends to optimize the density fluctuations. This peculiar electrostatic response is strictly of a quantum mechanical origin; in the classical limit, i.e. for low magnetic fields, there is no coupling between a magnetic field and the charge density at equilibrium. The quantum effects occur for high magnetic fields, of few Tesla or more, when the corresponding magnetic length is much smaller than the modulation period. [15]

The time-dependent density response function and the collective excitations of the 2DEG in the presence of a periodic electric or magnetic field have also been discussed recently by Stewart and Zhang. [16,17] They calculate the response function in the random-phase approximation finding subsingularities related to the transition energies near the van Hove singularities (vHS) of the energy bands. However, they neglected the influence of the Coulombian electron-electron interaction on the energy spectrum, and thus their results, even qualitatively correct, are restricted to a very weak modulation amplitude, much smaller than the cyclotron energy.

We have previously studied the FIR absorption by the magnetoplasmons of the 2DEG subjected to a short period unidirectional electric modulation and a uniform perpendicular magnetic field and discussed the effect of the modulation on the absorption peaks. [18] The peaks specific to the homogeneous system acquire an internal structure which can be related to the energy spectrum. The periodic Landau bands have vHS in the center and at the edges of the Brillouin zones. For a modulation period much longer than the magnetic length

the Landau bands are nearly parallel, with energy separation equal to the cyclotron energy. In this case the absorption spectrum is nearly identical to that in the homogeneous system. For a shorter modulation period the bands cease to be parallel, and the absorption peaks split into several sub-peaks following the dominant excitation energy which occur near the vHS. [16–18] In the present paper we want to discuss the absorption spectra in the presence of a magnetic modulation.

We compare the magnetic modulation effects with the situation of the homogeneous or electrically modulated 2DEG in the presence of a uniform and perpendicular (external) magnetic field. Therefore we will always keep a uniform field component, B_0. In the typical experimental realizations the period of the spatial magnetic modulation is of a few hundred nanometers. We will thus consider a period $a \gg l_0$, where $l_0 = (\hbar/eB_0)^{1/2}$. Within this restrictions we expect a dispersion of the excitation energies, determined by the cyclotron energy corresponding to the local magnetic field. In other words, the Landau bands will be non-parallel even for a weak magnetic modulation, and thus even a small nonuniform field component will have important effects. Therefore we expect FIR absorption spectra as rich in structure as for a short period electric modulation.

For the numerical calculations the material constants are chosen for GaAs: the effective mass $m_{eff} = 0.067m$, the dielectric constant $\kappa = 12.4$, and we assume spin degeneracy. The average electron density is fixed to $\rho_0 = 1.446\ 10^{11}\ \mathrm{cm}^{-2}$, such that $\nu B_0 = 6$ Tesla, where ν is the filling factor.

II. MODULATION MODEL AND SINGLE-PARTICLE STATES

We consider a 2DEG located in the plane $\mathbf{r} = (x, y)$, and an external magnetic field whose component along the z axis is independent of y, but periodic in the x direction, having the simple form

$$B(x) = B_0 + B_1 \cos Kx\,. \tag{2.1}$$

We assume the 2DEG to be ideally thin, such that it does not feel any in-plane component of the magnetic field. We describe the system with the Hamiltonian $H = H_0 + V_H$, where $H_0 = (\mathbf{p} + e\mathbf{A})^2/2m_{eff}$ is the noninteracting term, and V_H is the Hartree potential, self-consistent with the charge density.

In the Landau gauge, which is appropriate for our problem, the vector potential is

$$\mathbf{A} = (0, B_0 x + \frac{B_1}{K} \sin Kx)\,, \tag{2.2}$$

and with this choice the wave functions of the electrons have the form

$$\psi_{nX_0}(x, y) = L^{-1/2} e^{-iX_0 y/l_0^2} \phi_{nX_0}(x)\,, \tag{2.3}$$

where $n = 0, 1, ...$ and X_0 is the center coordinate. The y coordinate is thus isolated in a simple plane-wave, and ϕ_{nX_0} are the eigenfunctions of the reduced Hamiltonian

$$H_{0,x} = \hbar\omega_0 \left[-\frac{l_0^2}{2} \frac{d^2}{dx^2} + \frac{1}{2l_0^2} \left(x - X_0 + \frac{s}{K} \sin Kx \right)^2 \right]\,, \tag{2.4}$$

in which $s = B_1/B_0$ measures the modulation strength, and $\omega_0 = eB_0/m_{eff}$.

We will restrict ourselves to the case when $s < 1$, that is the magnetic field (2.1) vanishes nowhere. In this case we can expand the wave functions ϕ_{nX_0} in the basis of the Landau wave functions (i.e. the eigenfunctions of $H_{0,x}$ for $s = 0$) and we can neglect the contribution of very high Landau levels. Since we are interested in a magnetic modulation with a period $a = 2\pi/K$ much longer than l_0, $Kl_0 \ll 1$, we expect the perturbed wave functions ϕ_{nX_0} to preserve the localized character of the Landau wave functions, and only their "center of weight" eventually shifts from the center coordinate X_0 to a new position X_1. [15] The relation between X_1 and X_0 can be found by linearizing the sine of Eq.(2.4), $\sin Kx \approx \sin KX_1 + K(x - X_1)\cos KX_1$, and by reconstructing a new parabolic potential centered on X_1 one gets

$$H_{0,x} \approx \hbar\omega_0 \left[-\frac{l_0^2}{2}\frac{d^2}{dx^2} + \frac{(1 + s\cos KX_1)^2}{2l_0^2}(x - X_1)^2 \right] , \tag{2.5}$$

and

$$X_1 = X_0 - \frac{s}{K}\sin KX_1 \, . \tag{2.6}$$

The Hamiltonian (2.5) is the same as (2.4) with $s = 0$, but with modified cyclotron frequency $\tilde{\omega}_0 = \omega_0(1 + s\cos KX_1)$ and magnetic length $\tilde{l}_0 = l_0/\sqrt{1 + s\cos KX_1}$. The nonuniform magnetic field has thus also a scaling effect on the width of the wave functions, periodically in the center-coordinate space. The Landau bands resulting from Eq.(2.5) are thus

$$E_{nX_0} = \left(n + \frac{1}{2} \right) \hbar\tilde{\omega}_0 \, , \tag{2.7}$$

where one can directly see the spatial dispersion of the magnetic field, Eq.(2.1), and the Landau levels corresponding to the local cyclotron energy $e(B_0 + B_1 \cos Kx)/m_{eff}$, as we anticipated in the Introduction. The energy gap in the center and at the edges of the Brillouin zone, that is in the region where the density of states has vHS, is $\hbar\omega_\pm = \hbar e(B_0 \pm B_1)/m_{eff}$.

The electrostatic Hartree energy is determined by the charge distribution,

$$\rho(\mathbf{r}) = 2\sum_{nX_0} \mathcal{F}(E_{nX_0}) \mid \phi_{nX_0}(x) \mid^2 = \sum_{m \geq 0} \rho_m \cos mKx \, , \tag{2.8}$$

$\mathcal{F}$ being the Fermi function, and the Coulomb potential,

$$V_H(x) = \frac{e^2}{\kappa}\int d\mathbf{r}' \frac{\rho(\mathbf{r}')}{\mid \mathbf{r} - \mathbf{r}' \mid} = \frac{2\pi}{K}\frac{e^2}{\kappa}\sum_{m \geq 1}\frac{\rho_m}{m}\cos mKx \, . \tag{2.9}$$

The system is assumed electrically neutral, and thus the average density ρ_0 does not enter Eq.(2.9), but only determines the chemical potential contained in the Fermi function. The effect of the Hartree potential is a strong reduction of the density fluctuations imposed by the nonuniform magnetic field and also a reduction of the width of the Landau bands intersected by the Fermi level. [15]

In the numerical calculations we diagonalize iteratively the Hamiltonian $H_{0,x} + V_H(x)$ self-consistently with the particle density, by mixing up to 10 Landau levels, and using up

to 20 Fourier coefficients of the density. In Fig. 1 and Fig. 2 we show two characteristic energy spectra and density profiles for the self-consistent problem. We have fixed $B_1 = 0.2$ T. The modulation period is also fixed, $a = 500$ nm. We formally keep a finite temperature in all the calculations, $T = 1$ K. In Fig. 1 $B_0 = 2$ T ($\nu = 3$) such that the Fermi level is in the second Landau band, which is very flat. The dispersion of this energy band is canceled due to the electrostatic response. In Fig. 2 we have the opposite situation, in which, due to the low density of states at the Fermi level the energy dispersion is broad. But even if the Landau bands given by Eq. (2.7) may be strongly distorted due to the Hartree potential, the excitation energies $E_{n+1,X_0} - E_{n,X_0}$ are still given to a good approximation by $n\tilde{\omega}_0$. This is because the main action of the periodic magnetic field component, at least in the regime we are interested here, is on the degeneracy of the original Landau levels. Thus, for $Kl_0 \ll 1$ we may speak about a local character (in space) of the Landau-level degeneracy, as the mechanism of the static density response, and thus for the induced electrostatic potential.

The self-consistent electron-density profile may be smooth, like in Fig. 1(b), but also more complicated, with higher harmonics like in Fig. 2(b). In the latter case, in the spatial regions corresponding to the intervals of the Brillouin zone with an integer filling factor the density is constrained to follow proportionally the external magnetic field, Eq. (2.1), while in the regions with non-integer filling factors, the filling factor itself is adjusted by the electrostatic response, and the spatial variation of the density and of the magnetic field may be opposite. Another reason for the short-range oscillations of the static induced density is the combined effect of the shift of the wave functions, described by Eq. (2.6) and of the scaling of the wave function width from l_0 to $\tilde{l}_0$. Both these effects are periodic but out of phase. Consequently, short harmonics, with a characteristic length of a few l_0 are present in the static density response. Such oscillations are most prominent for not fully occupied Landau bands, Fig. 2(b). [15]

III. COLLECTIVE OSCILLATIONS AND ABSORPTION SPECTRA

We investigate the plasma oscillations and the energy absorption of the system with the help of the dielectric matrix $\varepsilon_{GG'}(q,\omega)$, G and G' being vectors in the reciprocal space, i.e. they belong to the set mK, $m = 0, \pm 1, \pm 2, \ldots$. We assume a weak time-dependent linearly polarized electric field incident on the system, of the form $\mathbf{E} = (E_x, 0)$, E_x having only one spatial Fourier component, of wave vector $q = (q_x, 0)$, and also only one temporal Fourier component, of frequency ω.

In the most common experimental situation the electric field is modulated by a system of grating couplers. [19,20] In our case such a modulation, with $q_x \equiv q = K$, can be produced by the magnetic or by the superconducting metallic strips themselves.

We calculate the dielectric matrix in the random-phase approximation, which is consistent with the Hartree approximation we adopted for the energy spectrum in the ground state, such that

$$\varepsilon_{GG'}(q,\omega) = \delta_{GG'} - \frac{2\pi e^2}{\kappa|q+G|}\chi_{GG'}(q,\omega)\,, \tag{3.1}$$

where the dielectric susceptibility $\chi_{GG'}$ is given by the Lindhard formula adapted to the periodic system, [24]

$$\chi_{GG'}(q,\omega) = \frac{1}{a\pi l^2}\sum_{nn'}\int_0^a dX_0$$
$$\times\frac{\mathcal{F}(E_{n,X_0}) - \mathcal{F}(E_{n',X_0})}{E_{n,X_0} - E_{n',X_0} - \hbar\omega - i\eta}$$
$$\times\,\mathcal{J}_{nn';X_0}(q+G)\mathcal{J}^*_{nn';X_0}(q+G')\,. \tag{3.2}$$

Here $\eta \to 0^+$ is the adiabatic switching parameter and $\mathcal{J}_{nn';X_0}(q) = \langle\psi_{n',X_0}|e^{iqx}|\psi_{n,X_0}\rangle$.

Due to the periodicity of the system $\varepsilon_{GG'}(q,\omega) = \varepsilon_{G-K,G'-K}(q+K,\omega)$. This property ensures the periodicity of the plasma frequencies, i.e. the zeros of $\det\varepsilon_{GG'}$ satisfy $\Omega(q) = \Omega(q+K)$, leading to a multiplication of the branches corresponding to the homogeneous system.

In the FIR spectroscopy however, only a few modes are observable, i.e. those with a non-vanishing oscillator strength. The absorbed power can be calculated from the Joule law of heating, expressible as [21]

$$P(q,\omega) = -\frac{\omega}{4\pi}\mathrm{Im}\,\varepsilon^{-1}_{GG}(q_1;\omega)\,q\,|\phi_{ext}(q,\omega)|^2\,, \tag{3.3}$$

where $q = q_1 + G$ with q_1 being the wave vector reduced to the first Brillouin zone, $0 \le q_1 \le K$, and $\phi_{ext}(q,\omega)$ denotes the electric potential of the incident field. For simplicity, we will normalize the external potential in Eq.(3.3) such that $\frac{1}{4\pi}q|\phi_{ext}(q,\omega)|^2 = 1$. In order to evaluate the absorbed power we will assume a certain dissipation in the system, by using a finite adiabatic switching parameter in Eq.(3.3), $\eta = \hbar\omega_0/50$.

In Fig. 3 we compare the absorption for the homogeneous system with that for a magnetically modulated system, for several wave vectors, with the same uniform magnetic field $B_0 = 2$ T. In the homogeneous system, for a low q, only the lowest magnetoplasma mode is active, according to the well known equation

$$\omega^2 = \omega_0^2 + 2\pi\rho_0 e^2 q/(\kappa m_{eff})\,. \tag{3.4}$$

With increasing q the oscillator strength of the higher mode, around $2\omega_0$, progressively increases, starting at about $q = 0.4K$. The high modes at $n\omega_0$, $n = 2, 3, ...$ are known as Bernstein modes. [22,23] In the modulated system we observe, in addition, intermediate modes, with $\omega_0 < \omega < 2\omega_0$, and some other small peaks outside this interval.

In Fig. 4(a) we show the dispersion of the first two modes for the non-periodic system versus the absolute wave vector q. In Fig. 4(b) we have folded the curves up to $q = 3.8K$ into the first Brillouin zone. Due to the reflection symmetry of the problem as imposed by the choice of the periodic magnetic field, Eq.(2.1), we have restricted Fig. 4(b) to half of the unit cell in the wave-vector space, $0 \le q_1 \le K/2$, the next half, corresponding to $K/2 \le q_1 \le K$, being the reflection of the first half. In the modulated system one expects a deviation of the frequency curves from those of Fig. 4(b) and a splitting at all the crossing points [24]. However, the peak structure of the traces in Fig. 3(b) can be completely understood with Fig. 4(b). For the modulated system several frequency branches (bands) become active for any value of q.

As we already mentioned, in the periodic system only the frequency dispersion is periodic with q, not the oscillator strength, which is the most relevant quantity in the absorption spectroscopy. For a low q the modulated system behaves indeed like the homogeneous one,

but with increasing q, the oscillator strength distributes to the neighouring branches. For instance, for $q = 0.4K$ and $q = 0.6K$, both corresponding to $q_1 = 0.4K$ in Fig. 4(b), instead of one peak for $1.2\omega_0 < \omega < 1.4\omega_0$, belonging to the lowest and respectively to the next frequency band shown in Fig. 4(b), we see in Fig. 3(b) both frequencies active. At $q = 0.8K$ we still have a peak on the lowest band at $\omega \approx 1.1\omega_0$, at the same position as for $q = 0.2K$, and also several peaks on the upper branches. The shoulder at about $\omega = 1.5\omega_0$ comes from the third band, and it becomes the main peak at $q = 1.2K$ ($q_1 = 0.2K$). The frequency of the main peak for $q = 0.8K$ is no more resolved for $q = 1.2K$, but the frequency of the lowest branch is still observable. We see further in each trace with $q > 0.8K$ two peaks, below and above $\omega = 1.6\omega_0$ respectively, with no dispersion in q. They correspond to the splitting of the uppermost branches, which in the homogeneous system are close to $\omega = 1.6\omega_0$. These branches are the most perturbed by the modulation: the lower is shifted downwards and the upper is strongly pushed into the frequency gap, as shown by the dashed lines in Fig. 4(b). These dispersionless modes are the transition modes from one to another Bernstein harmonic.

Similar dispersionless peaks are obtained for $\omega \approx 1.9\omega_0$, $\omega \approx 2.0\omega_0$, and $\omega \approx 2.1\omega_0$, from the branches with frequencies above $2\omega_0$. The perturbation of these branches due to the modulation is however stronger than on the lower branches, below $2\omega_0$, because the excitation spectrum is now wider, between $2\omega_-$ and $2\omega_+$, i.e. between $1.8\omega_0$ and $2.2\omega_0$. Not all the observable peaks gain in intensity with increasing wave vector, see for example the uppermost modes for $q = 1.0K$, or $q = 1.2K$. In the latter case one peak between two absorption maxima is completely missing, that is the corresponding mode is inhibited (very low oscillator strength).

It is worth to mention that for the modulated system, in the gap region the integral in Eq.(3.2) is singular due to the poles of the integrand between $2\omega_-$ and $2\omega_+$, such that a numerical calculation of the gap modes as zeros of determinant of the dielectric matrix would become very delicate. Our analysis, with a finite imaginary part, is thus much more efficient.

In Fig. 5 we show the evolution of the peak structure of the absorption with increasing modulation amplitude, for $B_0 = 1.5$ T, that is for the situation when the Fermi level is in a gap, and the Landau bands are wide. Here we have fixed the wave vector $q = K$. Like in Fig. 3, the modulation activates the modes in the vicinity of the frequency branches of the unmodulated system. All the frequencies are in fact changed by the modulation, but the most significant change is again of those close to the frequency gap. The modes originating at the gap edges are shifted into the gap.

The peaks which evolve directly from the peaks in the unmodulated system can still be distinguished by the structure of the time-dependent induced density, rather than by the frequency position, as we want to demonstrate in Fig. 6 for $B_1 = 0.2$ T. In the upper panel we show the induced density for two frequencies in the unmodulated system. For $\omega = 1.20\omega_0$ (not shown in Fig. 5(a)) there is no absorption peak, and the particle density fluctuates in time such that the time-dependent induced density is always symmetric. That is no dipolar motion occurs within the modulation unit cell, $0 < Kx < 2\pi$. On the contrary, for a peak frequency, like $\omega = 1.69\omega_0$, the induced density is antisymmetric, dipole active. For the modulated system we first display, in the lower panel, a symmetric induced density, for a frequency where no absorption takes place. The symmetry is now only approximate, but the

absorption is insignificant at this frequency. At $\omega = 1.64\omega_0$ the induced density reminds of the dipolar motion in the homogeneous system. It is antisymmetric and the charge motion in the center of the unit cell is spread over a large region. Due to the modulation some additional density fluctuations are present at the edges of the cell. For $\omega = 1.52\omega_0$ the dominant motion is not far from the cell center, between $KX = 1$ and $KX = 4$, while for $\omega = 1.77\omega_0$ the density oscillations are confined at the cell edges. We ascribe this behavior to the specific energy spectrum. The excitation frequencies split into ω_+ and ω_- corresponding to the vHS of the energy bands, and thus to the flat regions in the energy spectrum. Since the modulation period is much larger than the magnetic length, those magnetoplasma modes with frequencies slightly below and slightly above the frequency in the homogeneous system, here $\omega = 1.52\omega_0$ and $\omega = 1.77\omega_0$ respectively, are thus determined by local oscillations of the electron density around the center and around the edges of the unit cell respectively.

The local character of the modes which are activated by the modulation is better seen for the group around $2\omega_0$. Even in the previous group we already see some weak high harmonics in the induced density, for $\omega = 1.77\omega_0$. Such short range oscillations become very pronounced for the higher frequencies. Again the middle frequency, $\omega = 2.11\omega_0$, preserves the long range behavior of the peak mode in the unmodulated system. For the lower frequency, close to $2\omega_-$, the center of the unit cell is more active, this time with very pronounced short range harmonics, and for the higher frequency, close to $2\omega_+$, the motion at the cell edges becomes the most dipolar active.

IV. COMPARISON WITH AN ELECTRIC MODULATION

For a comparison of the effects of a magnetic modulation with those of a pure electric modulation we show in Fig. 7 results obtained for $B_1 = 0$, but with a periodic electrostatic potential $V \cos Kx$, again of period 500 nm, and of amplitude $V = 20$ meV. The uniform magnetic field is $B_0 = 1.5$ T, and $B_0 = 2$ T, and the absorption spectra are represented by the dashed and respectively the full lines of Fig. 7(a). The wave vector of the incident field is again fixed to $q = K$. In order to obtain significant deviations from the results in the homogeneous system, comparable with those in the presence of the periodic magnetic field, we had to increase the strength of the electric modulation, $V/\hbar\omega_0$, almost two orders of magnitude above the strength $s = B_1/B_0$ of the magnetic modulation. The peak structure of the absorption spectra in Fig. 7(a) is still less rich than what we have shown previously for the magnetic modulation.

The reason is the much stronger static response of the density to the periodic electrostatic potential than to the periodic magnetic field, compare Fig. 7(b) with Figs. 1(b) and 2(b). Consequently, the screening effects in the presence of an electric modulation are much stronger than in the presence of a magnetic modulation. The Landau bands corresponding to Fig. 7(a) have the same width as those in Fig. 1(a), for $B_0 = 2$ T, and Fig. 1(b), for $B_0 = 1.5$ T. The difference consists in weaker deviations of the excitation energies from multiples of $\hbar\omega_0$ than for the magnetic modulation, that is the Landau bands are now nearly parallel. Consequently we see in the absorption spectra of Fig. 7(a) less modes, i.e. those present in the homogeneous system for $q = K$, see Fig. 3(a) and Fig. 5(a), plus only one or two additional modes from the neighboring branches which are activated by the modulation, see Figs. 4(b) and 5(b). The original frequencies are significantly perturbed only for $B_0 = 1.5$

T, where we observe the tendency of the marginal modes to penetrate the frequency gap, respectively, the peak slightly above $\omega = 1.7\omega_0$ and the shoulder below $\omega = 2\omega_0$.

The electron density induced by the incident field, shown in Fig. 7(c), is long ranged for the lowest peaks, e.g. for $\omega = 1.71\omega_0$, resembling that in the unmodulated system. Higher harmonics are now present only for the small peak at $\omega = 2\omega_0$, indicating the onset of local charge oscillations. Much weaker high harmonics are also observable in the modes at the higher frequencies. Since in the case of the magnetic modulation the electron density has higher harmonics already in the ground state, Fig. 2(b), short range dynamic fluctuations are favored in the presence of the incident field, and thus the local modes, while for an electric modulation such modes are possible only for a large modulation amplitude, or for a short period, i.e. for a steep modulation potential. To summarize, the influence of a long period magnetic modulation on the FIR absorption in a 2DEG is qualitatively similar to that of a short period electric modulation of a comparable strength.

V. CONCLUSIONS

We have analyzed the magnetoplasma modes perpendicular to a one-dimensional magnetic modulation. We have discussed the influence of the Landau band structure, the density profile, and self-consistent electrostatic screening on the structure of the absorption spectra. We compare these effects to the analog results obtained in the presence of an electric modulation. In order to keep the interpretation as simple as possible, we have chosen the simplest situation, when the incident electric field is oriented perpendicular to the modulation, and is uniform along the modulation, that is $q_y = 0$. This situation is usually encountered in experiments on grating coupler devices. In the magnetically modulated systems the modulation of the incident, external electric field, with $q_x = K$, is possible by the metallic strips already incorporated in the device for producing the periodic magnetic field. In order to explain the results we have also considered the more general case with $q_x \neq K$.

Due to the weak and short range static density response of the 2DEG to a periodic magnetic field, as compared to the response to a periodic electrostatic field of the same period and strength, the screening effect on the Landau band structure is much weaker in the former than in the latter case. Therefore the structure of the absorption spectra for the magnetic modulation is much richer than for the electric modulation. For simplicity, we have limited our discussion to the situation when the modulation strength is sufficiently weak, such that the absorption spectra can be related to the collective modes occurring in the unmodulated system.

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FIGURES

FIG. 1. Energy bands (a) and electron density (b) for a magnetic modulation with $B_0 = 2$ T and $B_1 = 0.2$ T. The dashed line shows the Fermi level.

FIG. 2. Energy bands (a) and electron density (b) for a magnetic modulation with $B_0 = 1.5$ T and $B_1 = 0.2$ T.

FIG. 3. Absorption spectra for $B_0 = 2$ T in a homogeneous system (a) and in a modulated magnetic field with $B_1 = 0.2$ T (b) for various wave vectors of the incident electric field perpendicular to the modulation.

FIG. 4. Magnetoplasma dispersion of the first two modes, in the homogeneous system, for $B_0 = 2$ T (a) and the back-folded dispersions for $q \leq 3.8K$ (b) where q_1 is the wave vector reduced to the interval $[0, K]$; due to the reflection symmetry we have shown only the interval $[0, K/2]$. The dashed line shows the dispersion of the observable absorption maxima in the modulated system with $B_1 = 0.2$ T, for $q \leq 1.5K$.

FIG. 5. Various absorption spectra for a magnetic modulation with $B_0 = 1.5$ T, and with an increasing B_1, with $q = K$ (a) and the back-folded dispersions of the first two modes in the homogeneous system, for $q \leq 4.4K$ (b). (Note that the dispersion depends on B_0, being thus different from that of Fig. 4.)

FIG. 6. Induced electron density for $q = K$, for several frequencies shown by the numbers inside the plots, in units of ω_0, for the homogeneous system (a) and for the magnetic modulation with $B_0 = 1.5$ T and $B_1 = 0.2$ T (b).

FIG. 7. Results for a pure electric modulation of amplitude $V = 20$ meV ($B_1 = 0$). Absorption spectra for $q = K$, with $B_0 = 2$ T and $B_0 = 1.5$ T, full and respectively dashed lines (a), electron density for $B_0 = 1.5$ T (b), and a few induced densities for $B_0 = 1.5$ T, for the frequencies shown by the numbers (in units of ω_0) (c).

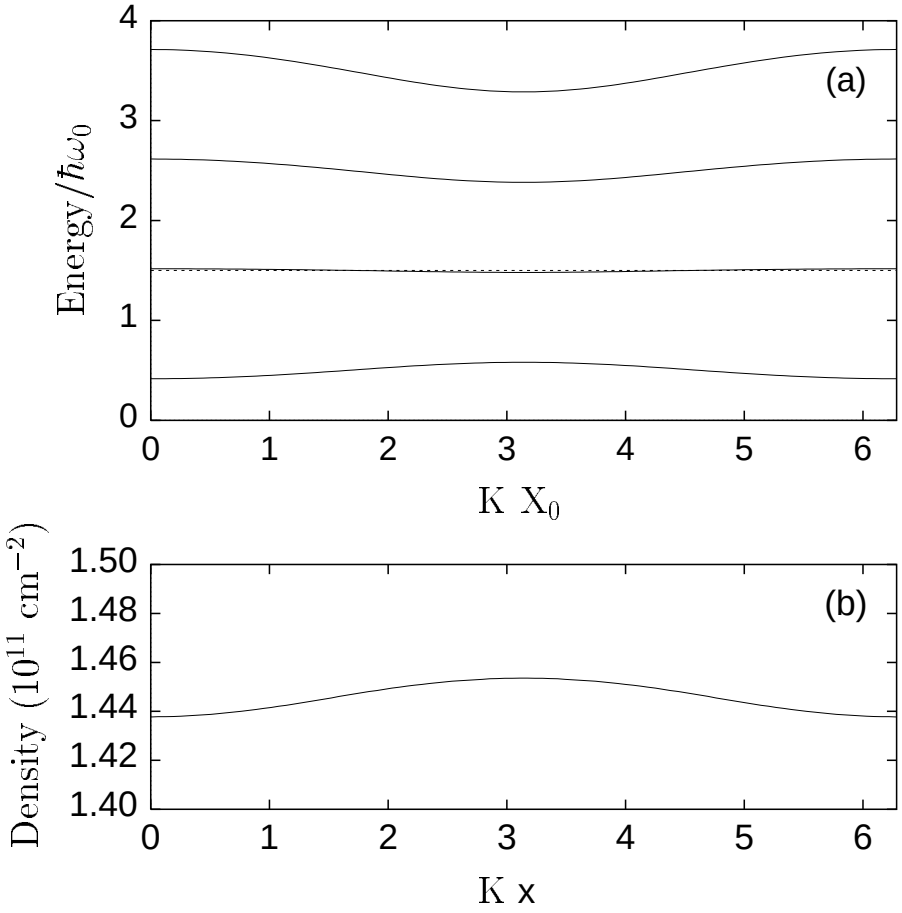

Fig.1

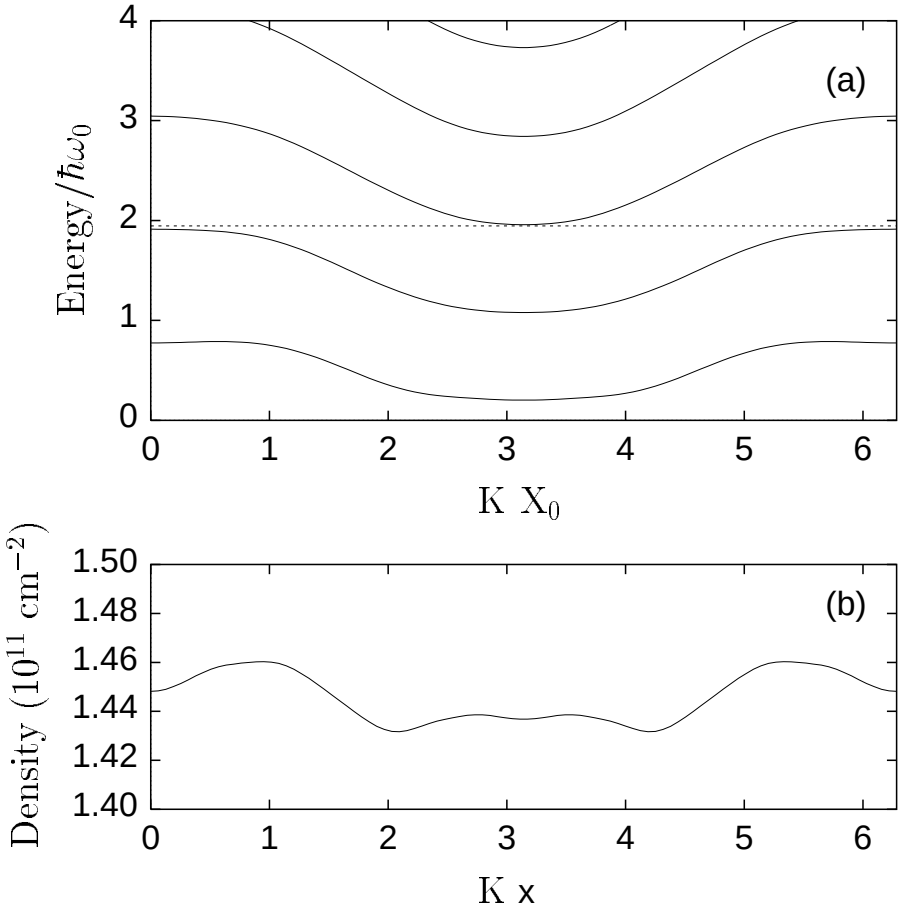

Fig.2

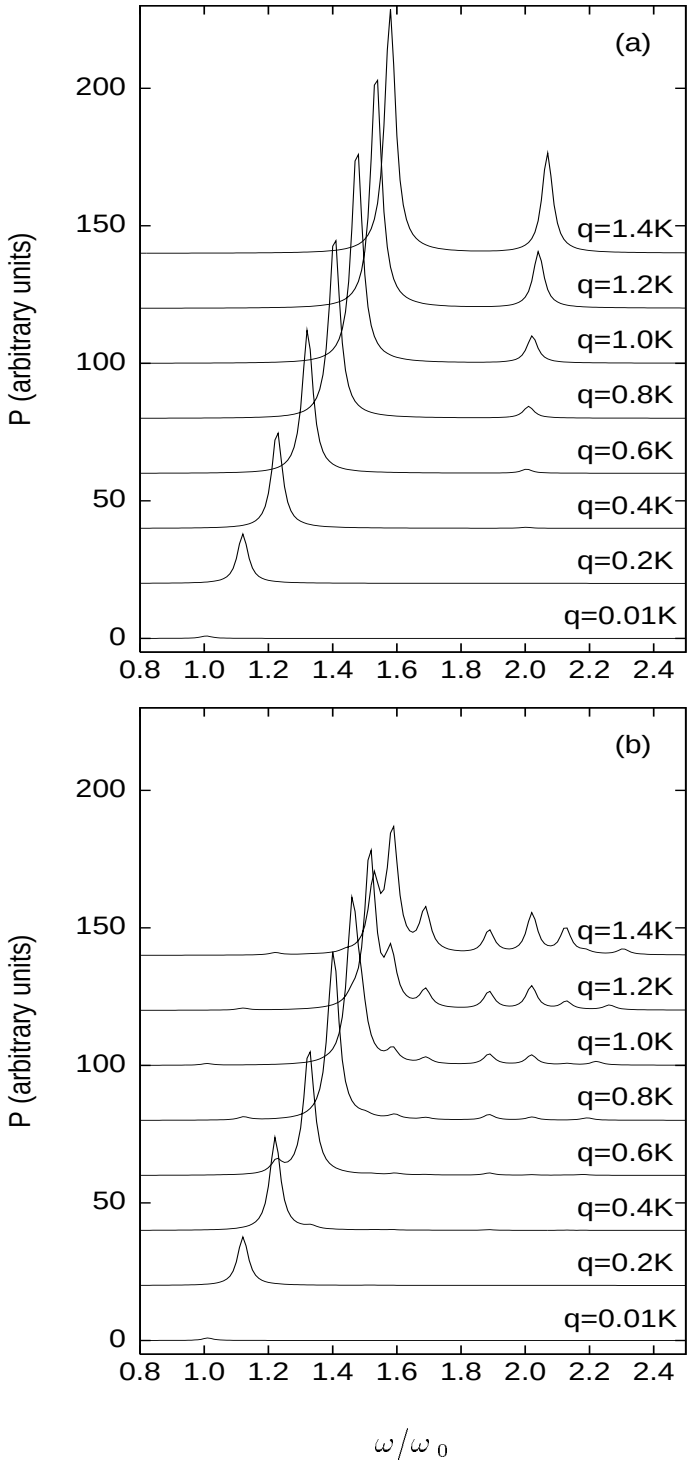
(a)
(b)
P (arbitrary units)
0
50
100
150
200
0.8
1.0
1.2
1.4
1.6
1.8
2.0
2.2
2.4
q=1.4K
q=1.2K
q=1.0K
q=0.8K
q=0.6K
q=0.4K
q=0.2K
q=0.01K
ω/ω 0

Fig.3

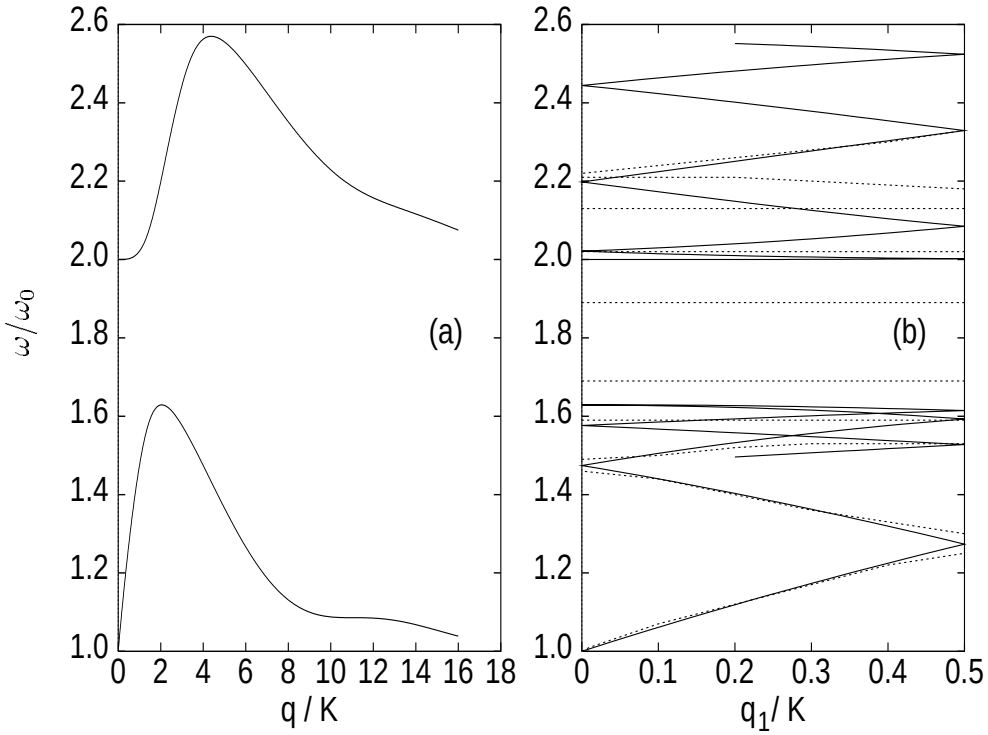
2.6
2.4
2.2
2.0
1.8
1.6
1.4
1.2
1.0
ω/ω0
(a)
(b)
0 2 4 6 8 10 12 14 16 18
q / K
0 0.1 0.2 0.3 0.4 0.5
q1/ K

Fig.4

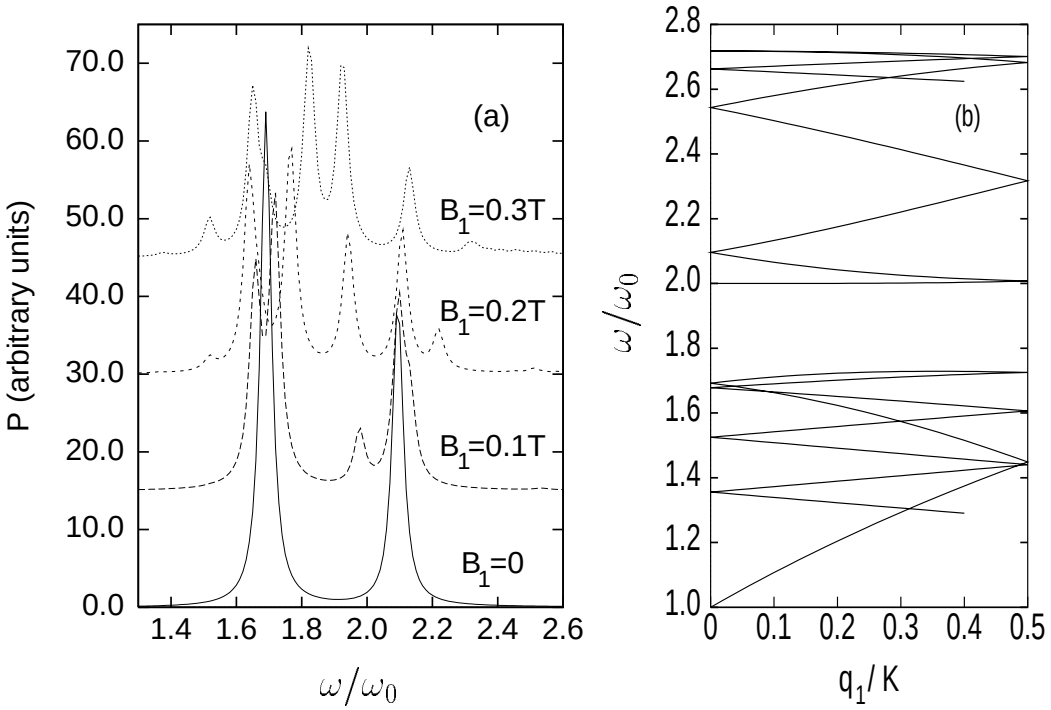

Fig.5

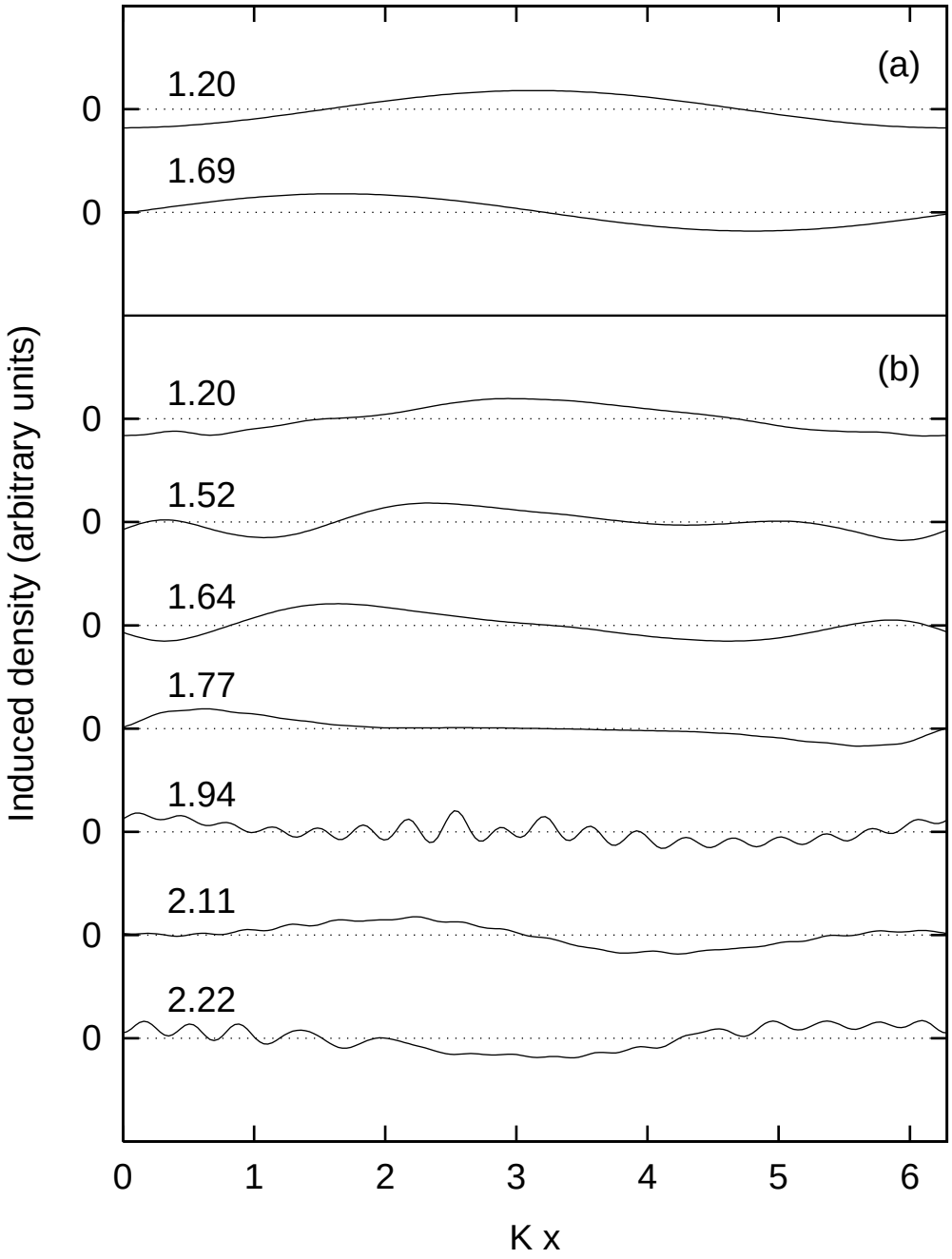
(a)
1.20
0
1.69
0
(b)
1.20
0
1.52
0
1.64
0
1.77
0
1.94
0
2.11
0
2.22
0
Induced density (arbitrary units)
0
1
2
3
4
5
6
K x

Fig.6

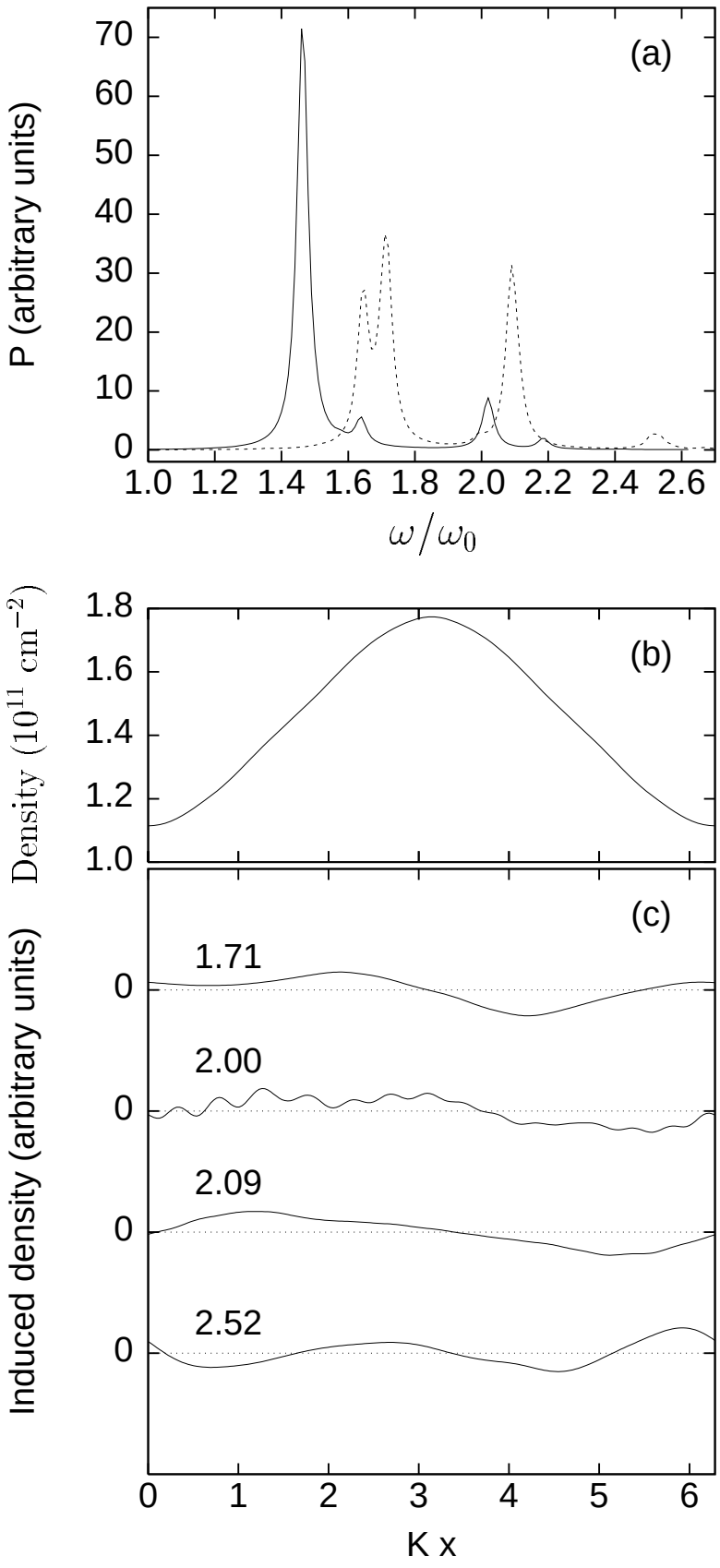
(a)
P (arbitrary units)
70
60
50
40
30
20
10
0
1.0 1.2 1.4 1.6 1.8 2.0 2.2 2.4 2.6
ω/ω_0
(b)
Density $(10^{11}$ cm$^{-2})$
1.8
1.6
1.4
1.2
1.0
(c)
Induced density (arbitrary units)
1.71
2.00
2.09
2.52
0
0 1 2 3 4 5 6
K x

Fig.7